

ANALYSIS OF THE SYNCHROTRON EMISSION FROM THE M87 JET

Sebastian Heinz¹

and

Mitchell C. Begelman²

JILA, University of Colorado and National Institute for Standards and Technology,
Boulder, Colorado 80309-0440³

Received _____; accepted _____

¹email address: heinzs@bogart.Colorado.edu

²email address: mitch@jila.Colorado.edu

³ also at Department of Astrophysical and Planetary Sciences, University of Colorado,
Boulder

ABSTRACT

We propose that the intensity changes and spectral evolution along the M87 jet can be explained by adiabatic changes to the particle momentum distribution function and the magnetic field. This is supported by the lack of any significant variation in the radio-to-optical spectral index along the jet and the moderate changes in radio brightness. Assuming a simple scaling law between magnetic field and density, we use the deprojection of a 2 cm VLA intensity map by Sparks, Biretta, & Macchetto (1996) to predict the spectral evolution along the jet.

We derive limits for the magnetic field and the total pressure by comparing our results with the spatially resolved fit to spectral data by Neumann, Meisenheimer, & Röser (1997) of a model spectrum that cuts off at $\approx 10^{15}$ Hz. To explain the weakness of synchrotron cooling along the jet, the magnetic field strength must lie below the equipartition value. Although the inferred pressure in the limit of nonrelativistic bulk flow lies far above the estimated pressure of the interstellar matter in the center of M87, bulk Lorentz factors Γ_{jet} in the range of $3 - 5$ and inclination angles $\theta_{\text{LOS}} \lesssim 25^\circ$ lead to pressure estimates close to the ISM pressure. The average best fit magnetic fields we derive fall in the range of $20 - 40 \mu\text{G}$, departing from equipartition by a factor $\approx 1.5 - 5$.

This model is consistent with the proposal by Bicknell & Begelman (1996) that the knots in the M87 jet are weak, oblique shocks. First-order Fermi acceleration will then have a minimal effect on the slope of the radio-to-optical spectrum while possibly accounting for the X-ray spectrum.

1. INTRODUCTION

The jet in M87 has been observed in many wavelength bands, yet some puzzles remain about the nature of its emissivity. Polarization observations in both the optical and radio have shown that the emission from the jet at wavelengths longward of $\approx 100 \text{ \AA}$ is most likely of synchrotron origin (Baade 1956, Owen et al. 1989). Collimated structure can be traced back to 0.01pc from the core using VLBI (Junor & Biretta 1995), and terminates $25''$ away from the core (2 kpc at the assumed distance of 17 Mpc) in the western radio lobe, with optical emission still detectable at this distance. The jet looks very similar at optical and radio wavelengths, although Sparks, Biretta, & Macchetto (1996, hereafter SBM) showed that differences (e.g., in the transverse brightness profiles) do exist.

Far from having a smooth appearance, the jet exhibits a series of bright knots at intervals roughly $2''.5$ apart. The nature of the knots is uncertain, but they are usually attributed to internal shocks from either flow instabilities (Bicknell & Begelman 1996, hereafter BB) or variable outflow at the source (Rees 1978). Both the knots and the interknot regions exhibit power law spectra of index $\alpha_R \approx 0.5$ in the radio, connecting to the optical data with a power law of index $\alpha_{RO} \sim 0.65$, and steepening to α_O between 1.2 and 1.8 in the optical. This steepening trend is also found in observations at infrared (e.g., Stocke, Rieke, & Lebofsky 1981) and ultraviolet (Perola & Tarenghi 1980) wavelengths, as confirmed recently by *HST* observations (Boksenberg et al. 1992, SBM). X-ray observations made with *Einstein Observatory* (Biretta, Stern, & Harris 1991) and *ROSAT* (Neumann et al. 1997a and Reynolds et al. 1996b) reveal X-ray emission from several spots along the jet (mainly the core and knot A, possibly also from knots D and B). However, the origin of the X-ray emission is unknown and it is not clear whether the spectrum breaks between optical and X-ray wavelengths to a spectral index of $\alpha_{OX} \approx 1.4$, with the X-ray emission still being of synchrotron origin, or whether the X-ray emission is produced by a different mechanism,

e.g., inverse Compton scattering or bremsstrahlung.

The onesidedness of the M87 jet has been interpreted as the result of relativistic beaming of the emission from an intrinsically bipolar jet, which implies bulk velocities corresponding to Lorentz factors of $\Gamma_{\text{jet}} \gtrsim 2$ with line of sight inclinations of $\theta_{\text{LOS}} \lesssim 30^\circ$. This would explain the existence of the radio lobes on both sides of the core. New infrared detections of counterjet emission (Stiavelli, Peletier, & Carollo 1997) seem to confirm this hypothesis. The Doppler beaming interpretation is bolstered by proper motion measurements of the knots (Biretta, Zhou, & Owen 1995), which show characteristic velocities $v_{\text{proper}} \approx 0.5 c$ in knots A and B, with some features exhibiting much larger proper velocities (a subfeature of knot D appears to show superluminal motion). These motions broadly support the interpretation of the knots as relatively weak, *oblique* shocks moving down the jet with pattern speeds significantly smaller than the bulk speed of the flow (BB).

One might hope to detect the effects of synchrotron cooling and relativistic particle acceleration by studying the spectrum as a function of position along the jet. Such measurements (SBM) show that the radio-to-optical spectral index, α_{RO} , is very nearly uniform, while the optical spectral index, α_{O} , is anti-correlated with the brightness, i.e., the optical spectrum is flatter in regions of higher intensity. Modeling the optical steepening as a high-energy cutoff imposed on a power law spectrum, Neumann, Meisenheimer, & Röser (1997, hereafter NMR) and Meisenheimer, Röser, & Schlötelburg (1996) find a corresponding correlation between brightness and cutoff frequency, i.e., a higher frequency at higher intensities. Both these results are striking in the lack of a strong secular decline in the cutoff frequency with distance from the core, as would be expected naively if the steepening were due to synchrotron cooling.

Indeed, these observational results do not compare well with simple quantitative models of synchrotron cooling in the M87 jet. The usual assumption of equipartition leads

to estimates of the magnetic field of order $300 \mu\text{G}$, with values up to $500 \mu\text{G}$ in knot A. The synchrotron lifetime for electrons with Lorentz factors of $\gamma \approx 10^6$, needed to produce the optical emission in a $300 \mu\text{G}$ field, is only 2.3×10^{10} sec, which, for mildly relativistic bulk velocities (e.g., $0.5 c$) implies a travel distance of less than 120 pc. But the projected length of the jet is about 2 kpc, and even the distances between the most prominent knots are longer than the estimated cooling length. Yet the spectrum between radio and optical bands remains remarkably constant along the jet, with only minor variations in the optical spectral index.

This presents a paradox: After a travel distance of 2 kpc one would expect the cutoff to have moved to below 10^{12} Hz, a factor of 10^3 smaller than what is actually observed. The discrepancy becomes worse if the magnetic field is stronger than the equipartition value, as suggested by Owen, Hardee, & Cornwell (1989) in order to explain the confinement of the overpressured jet via the magnetic tension force. (An alternate explanation for the confinement — that the radio cocoon surrounding the jet is overpressured with respect to its surroundings — has been proposed by BB.) Three explanations for the discrepancy have appeared in the literature:

1. First-order Fermi re-acceleration in the knots, interpreted as shocks, could produce high energy electrons (and possibly positrons) from a synchrotron-cooled distribution, with a power law of roughly the index observed. It could also explain the observed X-ray emission. However, since the power-law index produced by Fermi acceleration is a strong function of the compression ratio one would need some fine tuning to explain the observed constancy of the radio-to-optical spectral index, which does not equal the limiting value for a strong adiabatic shock. Furthermore, there does not seem to be a significant amount of cooling between the knots even at the highest optical frequencies observed, which would be expected for the assumed B -fields and interknot

distances.

2. Particles could be transported in a loss-free channel in the interior of the jet, with the bulk of the emission produced in a thin outer layer of high magnetic field strength (Owen et al. 1989). In this picture, the knots and filaments would be interpreted as instabilities with greatly increased magnetic field strength wrapped around the jet. The emission would then be fed by particles from this channel. The radio brightness profiles across the jet seem to suggest a limb-brightened emission, but a reinvestigation of the HST observations (SBM) shows that the optical emission is more concentrated to the inner regions of the jet. Also, SBM’s deconvolution places the brightest spots in the jet interior. This argues against a field-free zone in the jet interior.
3. On-the-spot reacceleration by a yet unknown process could maintain the cut-off particle momentum at the observed level, as has been proposed by Meisenheimer et al. (1996) in a model similar to ours (see §4.2). This has the advantage of explaining all the observed features, but invokes unknown physics to explain the apparent lack of cooling.

Inspired by the observed correlation between the emissivity variations and the cut-off frequency, by the newly deconvolved volume emissivity (SBM), and by the new evidence for relativistic bulk velocities (Stiavelli et al. 1997), we propose a simple way of explaining the observations. The only standard assumption we give up is the assumption of equipartition, which does not seem to have a very firm physical foundation anyway. Magnetic fields *smaller* than equipartition by a factor of $1.5 - 3$, coupled with bulk Lorentz factors in excess of $2 - 3$, can readily explain the lack of evidence for synchrotron cooling. In our model the fluctuations of the cutoff frequency are produced by weak shocks, so that the influence of the compressions on the plasma distribution function can be considered to be adiabatic. As a result, we are able to explain the general behavior of the cutoff reasonably

well; the magnetic fields we derive from our own fits to the data are below equipartition with a convincing level of confidence, but the inferred total pressures do not necessarily need to exceed the equipartition values. Additionally, for relativistic jets the equipartition values for B -field and pressure are less than in the nonrelativistic limit, so the pressures we derive can fall below the equipartition value in the nonrelativistic case.

We organize the paper as follows. In § 2 we present a nonrelativistic treatment of the synchrotron emissivity, taking into account cooling but assuming that the particles and fields respond adiabatically to changes in the flow density. Once relativistic effects are incorporated into the treatment in § 3, we use the data of SBM and NMR to constrain the magnetic field strength of the jet. Section 4 discusses confinement and stability of the jet in the light of the pressures derived from the results of § 3, the production of X-ray emission in knot A, polarization, and limits on the particle acceleration site; and § 5 gives a brief summary of the results and future prospects.

2. ADIABATIC EFFECTS ON SYNCHROTRON EMISSION

Our model rests on the hypothesis that Fermi acceleration is unnecessary to explain the fluctuations of radio-to-optical emissivity and cutoff frequency along the M87 jet. Given certain assumptions about the orientation and degree of disorder in the magnetic field, and the degree of anisotropy permitted in the relativistic electron distribution, we can relate changes in both the emissivity and the cutoff frequency uniquely to changes in the density of the jet fluid. These adiabatic effects are readily combined with the effects of synchrotron cooling (Coleman & Bicknell 1988). In effect, given the emissivity map of SBM, we can *predict* the run of the cutoff frequency along the jet, and vice-versa. Since we also have NMR’s observations of the cutoff frequency as a function of position, our adiabatic model is subject to a powerful self-consistency check.

Observationally, the main changes in emissivity and cutoff frequency are rather localized, and associated with the positions of the knots. These small-scale fluctuations therefore provide the strongest check on our assumption of adiabaticity. The *large-scale* trends then determine the best fit to the magnetic field strength. We have already seen that the apparent lack of a large-scale synchrotron cooling trend is incompatible with a magnetic field strength as large as the mean equipartition value (see Fig. 1). Neglecting relativistic and projection effects, it would require a magnetic field as low as $25 \mu\text{G}$ to obtain a cooling length of 2 kpc at optical frequencies. This is an order of magnitude smaller than the mean equipartition field and would require a total (particle + magnetic) pressure of $1.2 \times 10^{-7} \text{ dyn cm}^{-2}$ to produce the observed average amount of synchrotron emission, compared to an equipartition value of $\approx 4.0 \times 10^{-9} \text{ dyn cm}^{-2}$.

As noted earlier, the physical basis of equipartition is weak. Estimates of B -fields in radio hot spots based on the synchrotron cooling time indicate that in some cases equipartition might be correct up to a factor of ~ 2 (Meisenheimer et al. 1989), but the conditions in the jet might very well be different from those in the lobes. We are therefore free to consider the magnetic field strength to be a free parameter. Applying our adiabatic model to the observational data, we can derive an estimate for the magnetic field strength. In the case of the M87 jet, this estimate lies *below* equipartition, even when relativistic effects are taken into account (§ 3).

Our model for the evolution of the particle distribution function follows that of Coleman & Bicknell (1988). We assume that pitch angle scattering due to plasma micro-instabilities keeps the particles close to an isotropic distribution in the fluid rest frame. In the absence of cooling, this would imply that the relativistic electrons respond to compressions like a $\gamma_{\text{adiabatic}} = 4/3$ (i.e., ultrarelativistic) fluid, but this behavior will be modified by synchrotron cooling. Because the magnetic field is frozen into the plasma, its strength should change as

the plasma density fluctuates along the jet.

Depending on the orientation and the degree of disorder of the field, its variation will depend roughly on the density change to some power ζ : $B(r) \propto \varrho(r)^\zeta$, where ϱ is the proper particle density and r is the distance from the core. For a completely disordered magnetic field, $\zeta = \frac{2}{3}$, whereas for a homogeneous field the power depends on the orientation of the field with respect to the compression normal, with $\zeta = 1$ for an orthogonal orientation, $\zeta = 0$ for parallel orientation. Since the polarization of the jet is of order 10%-20% in the interknot regions, compared to the maximum polarization of 70% for a homogeneous field, it is likely that the magnetic field has a disordered component, so it is reasonable to assume ζ to be of order $\frac{2}{3}$ (but see §4.4). The ordered B -component is aligned with the jet axis almost everywhere except in the brightest knots, as can be seen from polarization measurements (Biretta & Meisenheimer 1993). We do not know its orientation with respect to the compression normal, because the orientations of the (presumably oblique: see BB) shocks are unknown. However, our results are not very sensitive to what the actual value of ζ is, as we will show later. We therefore make the simplifying assumption of a single exponent describing the field variations. Using the scaling relation for the synchrotron emissivity of a power law momentum distribution ($f(p) \, d^3p \propto p^{-a} \, d^3p \propto p^{-a+2} \, dp$, corresponding to a spectral index of $\alpha = \frac{a-3}{2}$) under adiabatic compression (e.g., Coleman & Bicknell 1988):

$$j \propto (B \sin \vartheta)^{1+\alpha} \varrho^{\frac{2\alpha+3}{3}} \nu^{-\alpha}, \quad (1)$$

where ϑ is the angle between magnetic field and line of sight, we can express the field relative to its value at $r_0 = 0''.5$, the (arbitrary) injection point at which we start the calculation, as a function of the emissivity ratio $j/j(r_0)$:

$$B = B(r_0)(j/j(r_0))^\xi \quad (2)$$

where $\xi \equiv [1 + \alpha + (2\alpha + 3)/(3\zeta)]^{-1}$. A necessary condition for this approach to be valid (in addition to the assumed isotropy and the absence of Fermi acceleration) is the assumed

steady-state injection of relativistic particles and fields by the central engine, which allows us to relate densities and fields at each r to the corresponding values at the injection point r_0 at a given instant of time.

Is it plausible to neglect Fermi acceleration in the shocks that comprise the knots? Except for knot A, the brightness changes along the jet are moderate. In knot A, the brightest feature, the emissivity changes by a factor of order 10 (SBM), which, if entirely due to a sudden compression of the plasma, can be produced by a proper compression ratio of $r \approx 2.7$ as measured in the respective rest frames of the plasma if we take $\zeta = \frac{2}{3}$. For the other knots we infer smaller density contrasts, $r \approx 1.5$. Consistent with this observation, we will henceforth take the knots to be weak shocks, in accordance with the suggestion by BB that the knots are highly oblique (and therefore weak) shocks. Thus, because Fermi acceleration leaves the spectral index unchanged if the shock is weak enough, we will henceforth neglect its effect on the cutoff frequency. Section 4 discusses Fermi acceleration in more detail, with particular focus on the possibility of Fermi acceleration occurring in knot A.

Furthermore, because the shocks are believed to be oblique, we take the fluid velocity to be constant to first order, both in magnitude and direction. For the shock jump conditions in the non-relativistic limit (which we consider in this section) the velocity component perpendicular to the shock plane v_\perp is inversely proportional to the density, thus for a proper compression ratio of 2.7, as seen in knot A, the perpendicular velocity component should change by a factor of 0.37. For highly oblique shocks, v_\perp is small compared to v_\parallel and the velocity will not change significantly. Moderate changes in velocity would be easy to incorporate in principle; yet with our current ignorance of the velocity field and shock parameters, such a level of detail is unwarranted. We will comment on the validity of this assumption in § 3. We also postpone a treatment of the motion of the knots until § 3, and

assume them to be stationary for the rest of this section.

With these assumptions we are ready to calculate the downstream cutoff frequency for a given initial cutoff momentum in the injected particle distribution and a given B -field at r_0 . We use the transport equation as presented by Coleman & Bicknell (1988):

$$\frac{df}{dt} + \frac{1}{3\varrho} \frac{d\varrho}{dt} p \frac{\partial f}{\partial p} = A p^{-2} \frac{\partial}{\partial p} (p^4 f) \quad (3)$$

written in the rest frame of the fluid. Here, $f = f(p)$ is the electron distribution function and $A = \frac{4e^4}{9m_e^4 c^6} B^2$ is the synchrotron loss term. The equation is valid for the assumed case of isotropy and negligible Compton losses (for a brief discussion of Compton losses see § 4.1). The solution of this equation is

$$f(p(r)) = f_0(p_0) \left(\frac{p_0}{p} \frac{\varrho^{\frac{1}{3}}}{\varrho_0^{\frac{1}{3}}} \right)^4 \quad (4)$$

where f_0 is the injected momentum distribution and

$$p(r) = \frac{(\varrho/\varrho_0)^{\frac{1}{3}} p_0}{(1 + p_0 \int_{t(r_0)}^{t(r)} A(t') (\varrho(t')/\varrho_0)^{\frac{1}{3}} dt')} \quad (5)$$

(see Coleman & Bicknell 1988) where the subscript 0 denotes the values at injection point r_0 .

Equation (5) describes how the momentum of a given particle changes along a streamline. Thus, if the distribution initially cuts off at $p_{c,0}$, we can calculate the cutoff momentum $p_c(r)$ downstream. Because in our model the density $\varrho(r)$ is proportional to $B(r)^{\frac{1}{\zeta}}$ and because we know the scaling of B with r from equation (2), we can eliminate ϱ and B/B_0 from equation (5). The remaining parameters are B_0 , $p_{c,0}$, and $j(r)/j_0$, the latter being provided by the Sparks et al. data.

The cutoff momentum p_c is related to the observed cutoff frequency ν_c by the expression

$$\nu_c = \frac{3e}{4\pi m_e^3 c^3} p_c^2 B \sin \vartheta. \quad (6)$$

Equation (6) contains another parameter, ϑ , the angle between the line of sight and the magnetic field. For now we shall set the factor $\sin \vartheta \approx 1$, which is valid for disordered fields, since 1) the regions in which the field is perpendicular to the line of sight have the highest emissivity, and 2) assuming randomly oriented fields, half of the field orientations lie in the range from 60° to 90° to the line of sight, i.e., $\sin \vartheta \geq 0.866$. Thus the cutoff frequency is mainly determined by field orientations close to 90° or $\sin \vartheta \approx 1$.

We can now determine the free parameters B_0 and $p_{c,0}$ by applying a least chi-squared method to fit the observed cutoff frequency $\nu_{c,\text{obs}}$ with the value determined from equations (5) and (6), using the emissivity map $j(r)/j_0$ provided by SBM. We prefer to average the emissivity across the jet, which minimizes small scale variations probably due to the deprojection procedure. Because NMR also averaged across the jet, this seems to be the most appropriate way of calculating the cutoff frequency. Figure 1 (calculated for a bulk Lorentz factor of $\Gamma_{\text{jet}} = 1.1$, a radio spectral index of $\alpha_R = 0.5$, $\theta_{\text{LOS}} = 90^\circ$, and $\zeta = \frac{2}{3}$) shows the observed cutoff frequency $\nu_{c,\text{obs}}$ (vertical bars) with error bars and the best fit curve (solid line), which seems to reproduce the scaling of the cutoff frequency reasonably well. (The radio spectral index seems to break to $\alpha \sim 0.65$ at ~ 10 GHz, so we have used both $\alpha = 0.5$ and 0.65 in our fits with insignificant differences in the average parameters but smaller chi-squared for $\alpha \sim 0.65$; see §3.) For comparison the plot also shows the best fit cutoff frequency for equipartition B -fields (dashed line). The mean B -field is of order $10 \mu\text{G}$, even smaller than the zeroth order estimate made at the beginning of this section. Assuming (arbitrarily) a lower cutoff at $\nu = 10^7$ Hz and the observed high-frequency cutoff at $\nu \approx 10^{15}$ Hz yields an average total pressure of $8 \times 10^{-8} \text{ dyn cm}^{-2}$ for the given parameters, compared to an equipartition value of $p \approx 3 \times 10^{-9} \text{ dyn cm}^{-2}$. In calculating absolute values for both pressure and B -field, projection (i.e., foreshortening and length scale) effects must be taken into account, since the emissivity was derived for a side-on view of the jet. This introduces a factor of $\sin \theta_{\text{LOS}}$ in intrinsic emissivity and pressure for a

given magnetic field and a factor of $(\sin \theta_{\text{LOS}})^{\frac{4}{7}}$ in equipartition pressure.

It seems that the proposed modest compressions can account for the fluctuations seen in the spectral cutoff, and the large scale decrease in ν_c is well reproduced. It would be helpful to determine both the emissivity and the spectral index maps from the same method and data, thus eliminating errors due to different reduction procedures. We comment on the deviations and uncertainties in this fit in § 3. It is important to note that this technique should be independent of what the actual shape of the particle distribution is, because it simply tracks the behavior of a single feature in the spectrum, which could be identified with either a break or a cutoff.

As seen in this section, one runs into problems with the jet pressure for nonrelativistic bulk velocities. Also, the observed mildly relativistic proper motion of the knots and the jet’s onesidedness favor a relativistic interpretation, as do the knot spacing and morphology (BB). The results indicate that Lorentz factors of order $2 - 5$ fit the observations best. In the next section we will investigate the effects of these suggested relativistic bulk velocities.

3. RELATIVISTIC EFFECTS

Relativistic motions not only explain the onesidedness of the M87 jet, but also help to solve the synchrotron cooling problem mentioned in the introduction. The travel time in the electron rest frame is reduced by a factor Γ_{jet} due to time dilation, the intrinsic emissivity is reduced by a factor $D^{2+\alpha}$ (where D is the Doppler factor $D = [\Gamma_{\text{jet}}(1 - \beta \cdot \cos \theta_{\text{LOS}})]^{-1}$), and the intrinsic cutoff frequency is Doppler shifted downward by a factor D . As a result, the apparent synchrotron lifetime can be a significant underestimate of the intrinsic value.

Biretta (1993) estimates the lower limit on the jet-to-counterjet radio brightness ratio to be $\geq 150 - 380$ (the higher value corresponds to the assumption that jet and counterjet

have identical appearance). Based on this limit we adopt line-of-sight angles $\leq 35^\circ$ and Lorentz factors ≥ 2 . Note that even Doppler factors smaller than unity can lead to a large jet-to-counterjet brightness ratio, as the counterjet brightness is severely reduced for large Γ_{jet} .

We repeat the analysis of § 2 using the emissivity profile of SBM, this time corrected for Doppler boosting and projection effects. Again, for fitting the cutoff frequency only emissivity ratios are important, so these corrections do not change the fitting procedure as long as changes in the bulk velocity can be neglected. We are confident that at least the direction of the flow is not changed significantly before knot C. The only knot for which such effects could be important is knot A, because it displays a jump in emissivity of 11, whereas in the other knots the emissivity is increased by a factor of order 3 only, which implies very moderate compressions. Using the relativistic continuity equation for an oblique shock we have estimated the post-knot A Lorentz factor to be $\Gamma_{\text{jet},A+} \sim 3$ for a pre-knot A $\Gamma_{\text{jet},A-} = 5$, a compression ratio of 3 and intrinsic obliquities $\sim 60^\circ$. Although the change in Γ_{jet} might seem large at first glance, the impact such a velocity change has on our fits is not large, as we will explain below. We modify equation (3) to follow the electron distribution in the *fluid rest frame* by replacing t with τ and $\frac{d}{dt}$ with $\frac{d}{d\tau}$, where τ is the proper time. The same changes apply to equation (5). Strictly speaking, the fluid frame is not an inertial frame and we would have to include accelerational terms into the equation, introducing an anisotropy. But our assumption should be adequate, provided that isotropization takes place over short enough scales.

Treating the response of the distribution function to compressions as adiabatic and assuming isotropy (i.e., an adiabatic index of $\frac{4}{3}$) we calculate the changes in B and cutoff momentum (measured in the fluid frame) from the emissivity changes in the fluid frame. The substitution $t \rightarrow \tau$ takes care of the time dilation effects.

In order to incorporate the observed motion of the knots, we must correct $A(\tau')$ and $\varrho(\tau')$ in equation (5) for light travel time effects between the source and the observer. This is because the knots move during the time it takes a particle to travel from r_0 to r . We need to know the ratios B/B_0 and ϱ/ϱ_0 experienced by a particle as a function of *proper time* τ' in order to be able to do the integration in equation (5). Because we infer the relative values of B and ϱ at a given position from the emissivity ratio $j(r_{\text{fluid}})/j_0$, it is important to know the velocity of the emissivity pattern.

We assume that the pattern of density and field fluctuations retains its shape and moves along the jet at a fixed speed v_{pattern} , taken to be smaller than v_{fluid} and set equal to $0.55\,c$ everywhere in our calculations for simplicity. If the present field distribution (i.e., at a given time $t = 0$ in our frame, corrected for light travel time effects) is expressed by $B(r)$, then the field at time $-t'$ and position r' is given by $B(r' + v_{\text{pattern}}t')$. Now, for a particle currently at r , the equation of motion is $r' = r - v_{\text{fluid}}t'$. Therefore, the field distribution experienced by the particle as a function of time is $B[r - (v_{\text{fluid}} - v_{\text{pattern}})t']$. Appropriate modifications to equation (5) are straightforward. The effect of the pattern speed on the result is not very dramatic, reducing χ_{min}^2 by about 6%.

With this set of assumptions we can once again proceed to integrate the modified equation (5) for various θ_{LOS} and Γ_{jet} . Using a minimum chi-square routine we can determine the best-fit values for B_0 , and p_0 . Relation (2) then yields $B(r)$.

Figure 2 shows a chi-square plot for $\Gamma_{\text{jet}} = 3$, $\theta_{\text{LOS}} = 25^\circ$, and $\zeta = \frac{2}{3}$. The equipartition value for the average B -field is shown as a shaded area at $89\,\mu\text{G}$. The upper limit on B_{mean} , set by $2\chi_{\text{min}}^2$ contours, lies at $49\,\mu\text{G}$, 75% above the best fit value, $B_{\text{mean}} = 28\,\mu\text{G}$. The lower limit set by $2\chi_{\text{min}}^2$ is $\approx 5\,\mu\text{G}$, 80% below the best fit value. The average equipartition field of $B_{\text{mean}} \approx 89\,\mu\text{G}$ lies above even the $5\chi_{\text{min}}^2$ contour. The lower limit on B is not nearly as strict, due to the fact that cooling is not dominant, i.e., we can produce a similar

spectral behavior by reducing the magnetic field and increasing the particle energy, which produces the tear-shaped appearance of the contours. Thus, strictly speaking, the best fit values for the B -field should be regarded as upper limits.

The reduced χ_{\min}^2 values (i.e., χ_{\min}^2 divided by the number of degrees of freedom) fall above 44, which is uncomfortably high. However, because we do not have formal errors for the emissivity deprojection by SBM, which will introduce a significant uncertainty, a high value for χ_{\min}^2 is not all that discouraging. Estimating the average uncertainty in the emissivity by comparing the averaged emissivity to that derived from taking only a slice along the jet yields an uncertainty of order 50%, which leads to uncertainties in the predicted cutoff frequency of roughly 20%. This is significantly higher than the formal error in NMR’s data and will reduce the χ_{\min}^2 by a factor of approximately 10.

The χ_{\min}^2 values are dominated by the region beyond knot A. The post-knot A residuals in our fit are not larger than the residuals in the pre-knot-A region, but because the post knot A region is brighter, the error bars on the measured cutoff frequency are smaller, which increases the χ_{\min}^2 . The deprojection procedure, which assumed an axially symmetric flow, breaks down beyond knot A, which will introduce significant uncertainty. Also, non-uniformities in the emissivity could lead to large errors if the optical emission peaks at different locations than does the radio emission. Field orientation effects and changes in Γ_{jet} and θ_{LOS} might also contribute to the error. We performed the same procedure just out to knot A and found that, with the *same* parameters, the reduced χ_{\min}^2 shrinks to 13. Leaving B and p_c as *free* parameters reduces χ_{\min}^2 to 10, but also reduces the B -field significantly. Because in this case the algorithm mainly fits the region around knot A (where the error bars are smallest), we cannot expect the global run of ν_c to have significant impact on the fit, which would be necessary to extract information about the average magnetic field. We conclude that the reproduction of fine detail is not satisfactory in the region beyond

knot A. However, the gross run of ν_c , which is principally responsible for constraining our parameters, is reasonably well reproduced.

The best-fit average magnetic field B_{mean} , as a function of Γ_{jet} and θ_{LOS} , is plotted in Figure 3. Figure 4 shows the average ratio $\langle B/B_{\text{eq}} \rangle$ as a function of Γ_{jet} for $\theta_{\text{LOS}} = 15^\circ$ to 30° in increments of 5° , and the area corresponding to the limiting jet-to-counterjet brightness ratio of 150 - 380. Note that the equipartition magnetic field has to be corrected by a factor of $D^{-\frac{2+\alpha}{3+\alpha}}(\sin \theta_{\text{LOS}})^{\frac{1}{3+\alpha}}$ for projection and Doppler boosting of the emissivity; this has already been taken into account in the figure. Clearly, for Γ_{jet} in the range 3 – 5 and $\theta_{\text{LOS}} \lesssim 25^\circ$ the departure from equipartition is not very large (roughly a factor of $0.2 < \langle B/B_{\text{eq}} \rangle < 0.6$).

In order to test the dependence of the best fit B_0 on the parameter ζ we have calculated the same curves for $\zeta = 1$ and $\zeta = \frac{1}{15}$. Figure 5 shows the fractional deviation $\left(\frac{\Delta B}{B}\right)_\zeta \equiv \frac{B_{\zeta_1} - B_{\zeta_2}}{B_{\zeta_1}}$ from the $\zeta_1 = \frac{2}{3}$ curve for models with $15^\circ \leq \theta_{\text{LOS}} \leq 30^\circ$. The deviation is small compared to the expected errors introduced by the simplifications we made and to the range in B allowed by our minimum chi-square procedure, at most 12% for $\zeta_2 = \frac{1}{15}$ and small Γ_{jet} . This is not a very reasonable value for ζ in any case, because the field has a random component, thus ζ should be higher, and the probability of the field being in the shock plane (thus having $\zeta = 1$) is twice as high as for the field being normal to the shock. We conclude that our ignorance of the precise behavior of B under compression is not a serious obstacle to the application of our model.

We also tested the impact a change in Γ_{jet} at knot A might have on our results. As we mentioned earlier, the best fit B -field values we derive are upper limits. This is the reason why a change in Γ_{jet} at knot A does not change our results significantly: generally, lower Lorentz factors require lower fields to explain the observed lack of cooling. If the jet is slowed down beyond knot A, we will need lower average fields to fit this

region. However, lowering the field does not change the quality of the fit much (the χ^2 is essentially unchanged), so the global field strength is simply set by the region with the lower Lorentz factor, $\Gamma_{jet,A+}$ (the relative scaling of B is still determined from equation 2, taking relativistic beaming into account). We have introduced by hand a change of Γ_{jet} at knot A into our model (we solved the continuity equation at the shock, assuming an obliquity of 60° , for the velocity change that would reproduce the observed emissivity jump of 11, including relativistic beaming and adiabatic compression), and calculated the fractional deviation $\left(\frac{\Delta B}{B}\right)_\Gamma \equiv \frac{|B_{\text{uniform}} - B_{\text{break}}|}{B_{\text{uniform}}}$ of the derived averaged B -field. Here, B_{uniform} is the best fit average B -field derived for uniform Γ_{jet} and B_{break} is the best fit field for a jet slowing down from $\Gamma_{jet,A-}$ to $\Gamma_{jet,A+}$ at knot A. Figure 5 shows $\left(\frac{\Delta B}{B}\right)_\Gamma$ for uniform jet models with Γ_{jet} set to either $\Gamma_{jet,A-}$ or $\Gamma_{jet,A+}$ (filled light and dark grey regions, respectively). The latter is always less than 18% for the parameter range we used. Note that for post-knot-A $\Gamma_{jet,A+}$'s above 5, the pre-knot-A $\Gamma_{jet,A-}$ exceeds 8.5, thus $\Gamma_{jet,A+} \gtrsim 7$ can be ruled out on the basis of gross energy balance arguments (see next section).

For completeness we have shown the deviation of the best fit average B -field $\left(\frac{\Delta B}{B}\right)_\alpha \equiv \frac{|B_{\alpha_1} - B_{\alpha_2}|}{B_{\alpha_1}}$ for a 2 cm radio spectral index of $\alpha_2 = 0.65$ instead of $\alpha_1 = 0.5$ as the black region in Figure 5. One can see that the difference is negligible compared to other uncertainties.

4. DISCUSSION

In the preceding sections we demonstrated that a) magnetic fields slightly below equipartition and b) moderately relativistic effects are able to explain the general behavior of the spectrum in M87. In this section we will examine the confinement properties of the jet and compare our model with a previous model by Meisenheimer et al. (1996). We also comment on the production of X-ray emission in knot A, and on the consistency of our

model with polarization measurements, and we estimate the minimum distance from the core at which particle acceleration has to occur.

4.1. *Confinement*

Naturally the question arises whether the jet can be confined under the conditions we proposed above. The usual assumption for a jet to be confined is that it is in pressure equilibrium with its surroundings. Alternatively, one could imagine the jet to be freely expanding into an underpressured surrounding medium.

BB argue that in order to produce shocks via Kelvin-Helmholtz instability, some interaction between jet and surrounding medium has to take place, as opposed to a free expansion scenario. They also show that the minimum Lorentz factor Γ_{jet} for a freely expanding jet with no cold matter content is at least 13, much higher than the values we have used above. Because we would most certainly fall out of the beaming cone for such a high Γ_{jet} jet, the intrinsic emissivity would be much higher than the observed value. As BB point out, the energy flux of the jet would far exceed the estimates made on the basis of the expanding bubble the M87 jet blows into the ISM. A Γ_{jet} that high would also raise questions about the location at which the jet is decelerated to nonrelativistic velocities, and seems inconsistent with the claimed detection of IR counterjet emission by Stievelli et al. (1997). We can therefore rule out the picture of a freely expanding jet. As a consequence we need a mechanism to provide confinement, i.e., we need to set the jet pressure in relation to the ambient pressure.

The ambient gas pressure in the center of M87 has been derived by White & Sarazin (1988) from fitting cooling flow models to the *Einstein* X-ray observations. The values they find fall into the range $p_{\text{ISM}} = 1 \times 10^{-10} \text{ dyn cm}^{-2}$ to $p_{\text{ISM}} = 4 \times 10^{-10} \text{ dyn cm}^{-2}$.

It is important to note that the pressure of the interstellar medium in M87 might not be representative of the pressure of the immediate environment of the jet. In fact, BB’s analysis of the helical Kelvin-Helmholtz instability leads to the conclusion that the ambient medium of the jet is significantly overpressured with respect to the interstellar medium in M87. Note also that the pressure in the knots might well exceed the ambient pressure without losing confinement, as long as the average pressure does not.

We have calculated the average total pressure p_{mean} in the jet from the averaged emissivity and the best-fit B -field for various angles and Lorentz factors, as shown in Figure 6. The assumptions we have made in constructing this plot are analogous to those of SBM, who (arbitrarily) assumed a lower cutoff at 10^7Hz , a high energy cutoff at 10^{15}Hz , a spectral index of $\alpha_{RO} = 0.5$, and equipartition between heavy-particle and electron energy. (Note: because the spectrum is steeper than $\alpha_R = 0.5$ above 10 GHz, our estimate of the pressure is likely to be an overestimate.) We have also assumed isotropic emission in the plasma rest frame by using an average value of $\vartheta = 54^\circ$ for the term $\sin^{1+\alpha} \vartheta$ in the emissivity equation (1).

For comparison we have also calculated the equipartition pressure and plotted the ratio of pressure to equipartition pressure in Figure 7. It is obvious that we are far above equipartition for small values of Γ_{jet} and large θ_{LOS} , but as we approach the favored range of $\Gamma_{\text{jet}} \geq 3$ and $\theta_{\text{LOS}} \lesssim 30^\circ$, p_{total} approaches the equipartition value. The exact value of the pressure depends critically on the details we put into the model spectrum. For a jet composed entirely of electrons and positrons, the pressure would go down by a factor of $\frac{1}{2}$, whereas the equipartition pressure would only decrease by a factor of $(\frac{1}{2})^{\frac{4}{7}} = 0.67$. The lack of information about the low-frequency spectrum inhibits any statements about the low-energy particle distribution. However, it is safe to assume that the power law does not continue down to non-relativistic energies.

It is obvious that a magnetic field far below equipartition alone cannot explain the behavior of the jet — it might account for the spectral changes but it requires the pressure to be much higher than that of the surrounding medium. Field orientations close to the line of sight will lead to an underestimate in emissivity, pressure, and intrinsic cutoff frequency and will require even lower magnetic fields and even higher pressures in order to prevent significant cooling. This changes as we increase Γ_{jet} : the inferred pressures are close to the value for the interstellar medium in M87 as derived by White & Sarazin (1988), for $\Gamma_{\text{jet}} \approx 3 - 5$ and $\theta_{\text{LOS}} \lesssim 25^\circ$. This, combined with the possible overpressure of the jet’s immediate environment relative to the ISM, leads to the conclusion that there is no confinement problem.

For small values of the magnetic field, one might ask if inverse Compton losses become dominant. A simple order of magnitude estimate shows that this is not the case. The ratio of synchrotron to inverse Compton loss timescales is equal to the ratio of photon energy density to magnetic field energy density (Rybicki & Lightman 1979). The magnetic field strength has been estimated above. To derive an estimate of the photon energy density produced by the synchrotron emission, we normalize to the radio luminosity corrected for beaming and projection effects and integrate over an $\alpha_{\text{RO}} = 0.65$ radio-to-optical power law that cuts off at 10^{15} Hz. This shows that the inverse Compton lifetime due to just the synchrotron radiation field of the jet is roughly an order of magnitude longer than the synchrotron lifetime for the parameter range we suggested above. The starlight background at the center of M87 also contributes to the photon energy density. Using an isothermal sphere profile, normalized to the total luminosity of M87, we arrive at a central photon energy density roughly an order of magnitude smaller than that of the magnetic field, small enough to justify the assumption of negligible Compton losses.

4.2. Comparison to Earlier Models

It is instructive to compare our model to an earlier *ad-hoc* model by Meisenheimer et al. (1996, see also Biretta & Meisenheimer 1993), which bears a lot of similarity to our model. They start from the same assumption that the spectral changes along the jet can be explained by simple compressions and assume that the cutoff momentum γ_c is almost constant along the jet, parameterizing it as a function only of the transverse jet radius (measured from the 2 cm radio map): $\gamma_c \propto R^{-\frac{1}{3}}$ — note that for an adiabatic compression of the plasma *transverse* to the flow γ_c goes as $R^{-\frac{2}{3}}$. They take the B -field to consist predominantly of a toroidal component, B_Φ , and hold the poloidal component B_z fixed. They determine the longitudinal compression ratio of B_Φ and of the particle density n from their fit to the cutoff frequency ν_c with equation 6. However, in an adiabatic compression, the cutoff momentum varies as $n^{\frac{1}{3}}$ and will therefore be affected by longitudinal compressions as well (here is where our assumption of a disordered field allows us to determine a relation between density and magnetic field, so we can solve equation 6 uniquely for B). They neglect the fact that the synchrotron emissivity is enhanced in adiabatic compressions by $n^{1+\frac{2}{3}\alpha} \cdot B^{1+\alpha}$ (equation 1) rather than $n \cdot B^{1+\alpha}$. Since the shocks might well be oblique, their assumptions that B_Φ scales as the longitudinal compression ratio and that B_z is constant might also not be valid.

Meisenheimer et al. (1996) favor an intrinsically onesided, subrelativistic jet, viewed close to perpendicular ($\theta_{LOS} \sim 90^\circ$). Knot A would be a head-on shock in this scenario. As we have mentioned above, for this set of parameters additional acceleration has to be provided to maintain the optical emission out to large distances from the core. Meisenheimer et al. (1996) favor an unknown global acceleration process to explain the constancy of the cutoff momentum. With these assumptions, their model yields similar results to ours in that it reproduces the small scale brightness variations on the basis of the changes in cutoff

frequency. Our model could thus be regarded as an extension of their approach, putting it on the theoretical basis of adiabatic expansions, with a different mechanism for providing the large scale constancy of the spectrum.

4.3. *X-ray Emission*

An important result from the analysis above is that cooling longward of UV energies is not important over the length of the jet — the proper time is reduced by a significant factor and the magnetic fields are small enough to leave the spectral shape unchanged. This conclusion begs the question of the origin of the X-ray emission detected by the *Einstein Observatory* and *ROSAT*. Both observations show emission from the core/knot D region and knot A.

4.3.1. *The Role of Particle Advection*

Ultra-high-energy particles, capable of radiating in the X-ray regime, could be carried out from knot D, where X-ray emission is observed, to knot A, and reaccelerated in the shock by the adiabatic compression mechanism discussed above. For this to happen the spectrum would have to break rather than cut off in the optical. The presence of a break instead of a cutoff would not change our fits, as long as the break is located above the frequency we fitted. We have calculated the behavior of particles with X-ray emitting energies along the jet and found that for our best fit B -fields cooling out to knot A will have produced a spectral cutoff at $\sim 10^{17}$ Hz — which is where most of the *Einstein* HRI’s sensitivity lies. Since the B -fields we derived are upper limits, a lower field could leave the distribution function unchanged even at such high energies. Therefore this mechanism of producing X-rays is marginally consistent with our model. It might also account for

at least part of the X-rays. Note that, as suggested by various authors (e.g., Biretta & Meisenheimer 1993), the X-ray emission could also be of non-synchrotron origin altogether.

4.3.2. *Fermi Acceleration at Knot A*

The compression of a factor ~ 3 inferred from the analysis above indicates that Fermi acceleration might be present in knot A, although the moderate change in emissivity and the constancy in radio and optical spectral index suggests that it might not be very efficient. It is possible that Fermi acceleration occurs at parts of the shock only, resulting in a particle distribution superposed of a compressed and a Fermi-accelerated preshock distribution.

In order for Fermi acceleration to take place at all the shock has to be subluminal, i.e., the intersection point of a given magnetic field line and the shock front has to move with a speed smaller than the speed of light, in which case we can find a frame in which the magnetic field is perpendicular to the shock front. This is the case for fields not too closely aligned with the shock plane. In the nonrelativistic case the field orientations leading to a superluminal shock are rare, and subluminal shocks are the rule rather than the exception, so one would expect Fermi acceleration to take place.

Because in relativistic shocks the percentage of superluminal field orientations rises sharply with Γ_{shock} , Fermi acceleration should become less important. In this limit, most of the particle acceleration would occur through the mechanism of “shock drift acceleration”. Begelman & Kirk (1990) presented a theory of this process valid in the relativistic case. They show that the adiabatic approximation is still accurate in the limit of $\Gamma_p \beta_p \lesssim 1$, where Γ_p is the upstream Lorentz factor in the perpendicular shock frame and β_p the corresponding velocity. We have calculated this quantity for various obliquities and field orientations appropriate for knot A and it seems to fall into the desired range. For a

disordered magnetic field we would have to average the resulting spectrum over all possible field orientations. The more superluminal the shock the less important would effects of Fermi acceleration be. Relativistic corrections to shock drift acceleration are only important for high Γ_p/β_p , which, again, depends on the field orientation. Averaging over all possible field orientations would probably render these corrections unimportant.

Fermi acceleration both changes the shape of an incoming power law spectrum and amplifies it. For an incoming electron spectrum of the form $f(p) = A_0 p^{-a}$ in the test particle limit (i.e., the pressure provided by the accelerated particles is negligible), and a nonrelativistic shock, the change of the spectrum depends on $s \equiv \frac{3r}{r-1}$, where r is the compression ratio. If $s < a$ the spectral index is changed to s . If $s > a$ the slope remains unchanged but the spectrum is still amplified by a factor $\frac{s}{s-a}$ (Kirk 1994). The radio spectral index is $\alpha \approx 0.5 - 0.65$, implying $a \approx 4 - 4.3$. To provide a boost in emissivity by a factor of ≈ 11 , s has to be $5.1 - 5.5$, implying a compression ratio of $\approx 2.2 - 2.5$ (assuming $\zeta = \frac{2}{3}$). Note that this is very close to the compression ratio one derives for an adiabatic compression. The spectral index produced by such a shock is $\alpha \approx 1 - 1.25$, consistent with the observed $\alpha_{\text{optical}} \approx 1.2$. Drury, Axford, & Summers (1982) showed that the produced power law softens as one departs from the test particle limit. Also, the simple treatment stated in this paragraph breaks down in the case of relativistic shocks, where the spectral index is no longer a simple function of the compression ratio.

Recent investigations by Ballard & Heavens (1992) have shown that oblique relativistic shocks can produce rather steep spectra, but other results indicate that they might be more efficient in accelerating particles (i.e., producing flatter spectra) than their nonrelativistic counterparts (Kirk & Heavens 1989). SBM find an optical-to-X-ray spectral index of $\alpha_{\text{OX}} \approx 1.4$ for the knot A region, which seems to be consistent with low-efficiency acceleration.

If Fermi acceleration were present at the shock and effective enough to change the spectral shape in the optical, it would no longer be feasible to use the data from the whole jet to determine the magnetic field. Rather, the same analysis could simply be carried out separately for the pre- and post-knot A regions of the jet. We would then have to make an estimate of the shock strength based on the known parameters in order to determine the ratio of pre- to post-shock Lorentz factor Γ . Even in this case we would need Lorentz factors of order $\Gamma_{\text{jet}} \approx 3$ to solve the cooling problem.

However, based on the observed mild spectral changes, and the assumption that pitch angle scattering is strong, we conclude that Fermi acceleration, if present, will not be efficient enough to affect the spectrum below the cutoff. Inefficient Fermi acceleration might very well be present in knot A, producing the X-ray emitting particles observed. Prediction of the produced high-energy spectral index has to wait for more conclusive results on Fermi acceleration at relativistic oblique shocks.

4.4. Polarization

An important complication to the treatment above is the fact that the magnetic field will not be completely disorganized — polarization measurements show that in some regions a homogeneous component is present. In fact, it is possible that cancellations between regions with homogeneous fields but different orientations occur along the line of sight (Meisenheimer 1992). In such a case, the assumption of disorganized fields, leading to $\zeta \approx \frac{2}{3}$, is no longer justified. However, since the impact that the parameter ζ has on the fit is minor (Fig. 5), we feel this caveat is not very severe and merely mention this complication here.

In addition to the unknown orientation of the jet itself, the field orientation is also

unknown. Since synchrotron emission depends on the magnetic field orientation to the line of sight ϑ as $\sin^{1+\alpha} \vartheta$, it can be strongly peaked away from the field direction. The cutoff frequency also depends on $\sin \vartheta$. Because the cutoff frequency is determined by sampling all regions along the line of sight, it is not obvious which value to choose for ϑ . Fortunately, in a domain of disordered field the regions with the field oriented close to perpendicular will contribute most of the flux, thus the error we make by setting $\sin \vartheta$ to 1 will not be too large. Note that we used a value of $\vartheta = 53^\circ$ in calculating the pressure, the appropriate average of $\sin^{1.5} \vartheta$ over 4π sterad.

Knot A shows a polarization of order 35% and a field orientation close to the shock plane. Neglecting relativistic effects, the small amount of upstream polarization, and the fact that shear might reduce the compression of the field, we can use the approximate formula (Hughes & Miller 1995)

$$\Pi \approx \frac{\alpha + 1}{\alpha + \frac{5}{3}} \cdot \frac{(1 - r^{-2}) \cos^2 \epsilon}{2 - (1 - r^{-2}) \cos^2 \epsilon}, \quad (7)$$

where Π is the fractional polarization, r is the proper compression ratio, and ϵ is the line of sight angle from the shock plane, to obtain a lower limit on the compression ratio in knot A of $r \approx 2$, valid for the inferred range of viewing angles with respect to the shock plane (BB). This is consistent with the compression ratios inferred above. Equation (7) also provides an upper limit of 35° on ϵ ; however, the shock is assumed to be oblique, so we cannot use this result to constrain θ_{LOS} .

4.5. Particle Acceleration Radius

Having an estimate of both magnetic field and cutoff momentum at the injection radius r_0 , we can now try to determine where the actual particle acceleration has to occur. We assume some radial dependence for the magnetic field in the *inner* portion of the jet

(i.e., smaller than $0.5''$), for example a power law: $B \propto r^{-\sigma}$. Furthermore we make the simplifying assumption of a constant Γ_{jet} . By also taking B to be proportional to ϱ^ζ , which determines the radial dependence of ϱ , we can invert equation (5) and solve for the radius at which the cutoff momentum approaches infinity, in other words, the minimum radius inside of which acceleration has to occur:

$$r_{\text{acc}} = r_0 \left(1 + \frac{1}{A \cdot r_0} \right)^{-\lambda} \quad (8)$$

where $A = \frac{4p_0 B_0^2 e^4 \lambda}{9\Gamma_{\text{jet}} c^7 m_e^4}$, and $\lambda \equiv \frac{1}{(2 + \frac{1}{3\zeta})\sigma - 1}$. The acceleration radius r_{acc} approaches zero if $A \cdot r_0 \rightarrow -1$. It is obvious that the estimate of r_{acc} depends critically on the value of σ . In Figure 8 we have plotted r_{acc} as a function of σ for $\theta_{\text{LOS}} = 25^\circ$ and $\Gamma_{\text{jet}} = 3$. In the same figure we have also plotted r_{acc} for the case in which B no longer scales like ϱ^ζ — in this case we have taken $\varrho \propto r^{-2}$ and used the same values for θ_{LOS} and Γ_{jet} (dashed line).

If the jet expands at constant opening angle and with uniform Γ_{jet} , the decline of B with radius should correspond to $\sigma \leq 2$, since the density scales like r^{-2} and dissipative effects will probably limit the rate of decline. Adopting an upper limit on the magnetic field at 0.01 pc of $B \leq 0.1$ G (Reynolds et al. 1996a) limits σ to values smaller than 1. For $\Gamma_{\text{jet}} = 3$, $\sigma = 1$, and $\zeta = \frac{2}{3}$, this constrains the acceleration radius to be $r_{\text{acc}} \leq 10$ pc, or $0''.06$. On the other hand, the radius of acceleration cannot lie inside the Schwarzschild radius of the central black hole, which is of order 10^{-4} pc for a $10^9 M_\odot$ black hole.

The dependence of r_{acc} on σ , ζ , and the core magnetic field is too strong to make any detailed predictions about where the acceleration actually has to occur. However, the r_{acc} -curve is rather flat throughout most of the possible range for σ , which suggests that the most plausible value for r_{acc} falls between 1 and 10 pc. This is intriguingly far away from the central engine.

5. CONCLUSION

We have proposed that the apparent lack of synchrotron cooling in the M87 jet most likely indicates the presence of a sub-equipartition magnetic field. While the total (particle + magnetic) pressure needed to explain the observed synchrotron emissivity is uncomfortably high in the nonrelativistic limit, Doppler beaming effects consistent with bulk Lorentz factors in the range 2 – 5 lower the pressure requirements considerably. Fluctuations in the synchrotron emissivity and spectral cutoff frequency are consistent with adiabatic changes in the magnetic field strength and particle energies that accompany compressions and rarefactions along the flow — Fermi acceleration along the flow is not necessary to explain the observations.

The knots are identified with relatively weak shocks, as inferred from other data by BB. The first-order Fermi acceleration expected to occur at such shocks, if any, would generate a particle energy distribution steeper than the $n(E) \propto E^{-2}$ needed to produce the radio-to-optical synchrotron spectrum. Thus, effects of particle acceleration along the jet might be apparent only shortward of the cutoff frequency, e.g., in the X-ray band. However, a disordered magnetic field or a “superluminal” field orientation with respect to the shock front (Begelman & Kirk 1990) could further reduce the efficiency of Fermi acceleration, hence we should continue to regard the origin of the X-ray emission as unknown.

The pressure estimates we derive are consistent with the assumption that the M87 jet is embedded in a moderately overpressured bubble, as suggested by BB. As a result, it seems that the set of parameters we have suggested above can solve the cooling problem of the M87 jet, as sub-equipartition fields are able to explain both the behavior of the cutoff frequency and the confinement of the jet.

Using derived values for the magnetic field and the cutoff momentum at r_0 we can put an upper limit of 10 pc on the radius at which most of the particle acceleration occurs.

Explaining why the acceleration is confined to a particular scale (which may be quite large compared to the size of the central black hole) poses an interesting problem for future work.

The method we have developed should be applicable to other radio sources as well, as long as propagation effects on the particle momentum distribution can be treated as adiabatic. Thus, as soon as comparable data becomes available for other jets, it would be easy to apply the same analytic techniques to them and obtain limits on the magnetic field and the pressure in these sources.

The analysis of the M87 jet we carried out in this paper could be improved by better deprojection models of the jet. Higher quality spectral data both in the near infrared and the X-ray would help to verify the assumption of constancy of the underlying spectral shape and decrease the uncertainty in the location of spectral features (the knowledge of which is essential for this technique to work). *HST/NICMOS* will be ideally suited to mapping out the jet properties in the near infrared, a spectral region in which the run of the spectral index is still fairly poorly determined. *AXAF*, providing imaging spectroscopy at sub-arcsecond resolution, will be the ideal tool to obtain more detailed information about the X-ray spectrum and emissivity along the jet, hopefully allowing us to determine the radiation mechanism at high energies.

The model we have described above fits in nicely with the new picture of M87 that has emerged in several recent papers. A moderately relativistic jet and magnetic fields somewhat below equipartition can explain the lack of cooling observed at optical wavelengths. Oblique shocks, weak enough to render Fermi acceleration unimportant, can explain both the small scale variability of the cutoff frequency and the brightness variations via adiabatic compressions. We believe this model can tie together most of the observations available today and will take us closer to the true nature of the jet.

We are grateful to Klaus Meisenheimer for an exceptionally thorough and detailed referee report, which helped us to improve the paper. Klaus Meisenheimer and Bill Sparks also provided us with the observational data files that we used to fit our models. This work has been partially supported by NSF grants AST-95 29175 and AST-91 20599 and a Fulbright UP95/96/97 grant.

REFERENCES

- Baade, W. 1956, ApJ, 119, 221
- Ballard, K.R., & Heavens, A.F. 1992, MNRAS, 259, 89
- Bicknell, G.V., & Begelman, M.C. 1996, ApJ, 467, 597 (BB)
- Begelman, M.C., & Kirk, J.G. 1990, ApJ, 353, 66
- Biretta, J.A. 1993, in Astrophysical Jets. Space Telescope Science Institute Symposium Series Vol.6, D. Burgarella, M. Livio, & C. O’Dea (eds.) (Cambridge: Cambridge University Press), 263
- Biretta, J.A., & Meisenheimer, K. 1993, in Jets in Extragalactic Radio Sources. Lecture Notes in Physics Vol. 421, H.–J. Röser & K. Meisenheimer (eds.) (Berlin: Springer), 159
- Biretta, J.A., Stern, C.P., & Harris, D.E. 1991, AJ, 101, 1632
- Biretta, J.A., Zhou, F., & Owen, F.N. 1995, ApJ, 447, 582
- Boksenberg et al. 1992, A&A, 261, 393
- Coleman, C.S., & Bicknell, G.V. 1988, MNRAS, 230, 497
- Drury, L.O., Axford, W.I., & Summers, D. 1982, MNRAS, 198, 833
- Hughes, P.A., & Miller, L. 1995, in Beams and Jets in Astrophysics. Cambridge Astrophysics Series Vol. 19, P.A. Hughes (ed.) (Cambridge: Cambridge University Press), 1
- Junor, W., & Biretta 1995, J.A., AJ, 109, 500
- Kirk, J.G. 1994, in Plasma Astrophysics. Saas-Fee Advanced Course Vol. 24, A.O. Benz & T.J.-L. Courvoisier (eds.) (Berlin: Springer), 225

- Kirk, J.G., & Heavens, A.F. 1989, MNRAS, 239, 995
- Meisenheimer, K., Röser, H.–J. & Schlötelburg, M. 1996, A&A, 307, 61
- Meisenheimer, K., Röser, H.–J., Hiltner, P.R., Yates, M.G., Longair, M.S., Chini, R., & Perley, R.A. 1989, A&A, 219, 63
- Meisenheimer, K. 1992, in Physics of Active Galactic Nuclei, W.J. Duschl & S.J. Wagner (eds.) (Berlin:Springer), 525
- Neumann, M., Meisenheimer, K., Röser, H.–J., & Fink, H.H. 1997a, A&A, 318, 383
- Neumann, M., Meisenheimer, K., & Röser, H.–J. 1997b, A&A, in preparation (NMR)
- Owen, F.N., Hardee, P.E., & Cornwell, T.J. 1989, ApJ, 340, 698
- Perola, G.C., & Tarengi, M. 1980, ApJ, 240, 2
- Rees, M.J. 1978, MNRAS, 184, 61P
- Reynolds, C.S., Fabian, A.C., Celotti, A., & Rees, M.J. 1996a, MNRAS, 283, 873
- Reynolds, C.S., Di Matteo, T., Fabian, A.C., Hwang, U., Canizares, C.R. 1996b, MNRAS, 283, L111
- Rybicki, G.B., & Lightman, A.P. 1979, Radiative Processes in Astrophysics, (New York: Wiley)
- Sparks, W., Biretta, J.A., & Macchetto, F. 1996, ApJ, 473, 254 (SBM)
- Stiavelli, M., Peletier, R., & Carollo, C.M. 1997, MNRAS, 285, 181
- Stoeckle, J.T., Rieke, G.H., & Lebofsky, M.J. 1981, Nature, 294, 5839, 319
- White, R.E., & Sarazin, C.L. 1988, ApJ, 335, 688

Fig. 1.— Measured cutoff frequency along the jet from NMR (vertical lines with error bars). Solid line: best fit curve as calculated from the emissivity measured by SBM for $\Gamma_{\text{jet}} = 1.1$, $\zeta = \frac{2}{3}$, and $\theta_{\text{LOS}} = 90^\circ$; dashed line: best fit curve for the case of equipartition calculated from SBM data for the same set of parameters.

Fig. 2.— Contour plot of the χ^2 values as a function of B_{mean} and the injection cutoff momentum P_0 for $\theta = 25^\circ$ and $\Gamma_{\text{jet}} = 3.0$. Shown are the contours corresponding to a reduced χ^2 of 45 and integer multiples of this value. The shaded line on the right of the plot shows the equipartition B -field for this set of parameters.

Fig. 3.— B -field averaged along the jet between $0''.5$ and $22''$ as a function of Γ_{jet} for $\theta_{\text{LOS}} = 15^\circ$ (solid line), $\theta_{\text{LOS}} = 20^\circ$ (dotted line), $\theta_{\text{LOS}} = 25^\circ$ (dashed line), and $\theta_{\text{LOS}} = 30^\circ$ (dashed-dotted). The grey region indicates the jet-to-counterjet brightness ratio limit of 150 (dashed grey boundary) - 380 (solid grey boundary, Biretta 1993).

Fig. 4.— Ratio of the best-fit B -field to the respective equipartition B -field B_{eq} averaged between $0''.5$ and $22''$ for $\zeta = \frac{2}{3}$ as a function of Γ_{jet} . The different values of θ_{LOS} are labeled similarly to Fig. (3). The grey region shows the jet-to-counterjet brightness limit (dashed grey boundary: 150, solid grey boundary: 380).

Fig. 5.— Fractional deviation $\frac{\Delta B}{B}$ of the best fit B -field (averaged between $0''.5$ and $22''$ for *a.* two different values of ζ : $\left(\frac{\Delta B}{B}\right)_\zeta \equiv \frac{|B_{\zeta_1} - B_{\zeta_2}|}{B_{\zeta_1}}$, where $\zeta_1 = \frac{2}{3}$ and either $\zeta_2 = \frac{1}{15}$ (hatched, solid boundary) or $\zeta_2 = 1$ (hatched, dotted boundary); *b.* a uniform jet, compared to a jet slowing down at knot A: $\left(\frac{\Delta B}{B}\right)_\Gamma \equiv \frac{|B_{\text{uniform}} - B_{\text{break}}|}{B_{\text{uniform}}}$, where B_{break} is the best fit average B -field for a jet slowing down from $\Gamma_{\text{jet},A-}$ to $\Gamma_{\text{jet},A+}$ at knot A, and B_{uniform} is the best fit average field for a uniform jet with $\Gamma_{\text{jet}} = \Gamma_{\text{jet},A-}$ (dark grey region, short dashed boundary) and $\Gamma_{\text{jet}} = \Gamma_{\text{jet},A+}$ (light grey region, long dashed boundary). The latter is plotted versus $\Gamma_{\text{jet},A+}$; *c.* two different 2cm radio spectral indices: $\left(\frac{\Delta B}{B}\right)_\alpha \equiv \frac{|B_{\alpha_1} - B_{\alpha_2}|}{B_{\alpha_1}}$ for our standard

value $\alpha_1 = 0.5$ and $\alpha_2 = 0.65$, shown as the black area. (The width of the band corresponds to θ_{LOS} between 15° and 30° in each case.)

Fig. 6.— Total pressure p_{mean} averaged along the jet between $0''.5$ and $22''$ as a function of Γ_{jet} . Labels according to Fig. (3). For comparison, the dashed-triple-dotted line shows the equipartition pressure for a non-relativistic jet seen edge on (i.e., $\theta_{LOS} = 90^\circ$). The hatched area shows the estimated ISM pressure (White & Sarazin 1988) in M87.

Fig. 7.— The ratio of the best-fit particle pressure to the respective equipartition value averaged along the jet as a function of Γ_{jet} . Labels according to Fig. (3).

Fig. 8.— The minimum acceleration radius r_{acc} , as a function of σ ($B \propto r^\sigma$) for $\Gamma_{\text{jet}} = 3$ and $\theta_{LOS} = 25^\circ$ in the case of $B \propto \varrho^\zeta$ (solid line). The dashed line shows r_{acc} for the case of $\varrho \propto r^{-2}$ and the same values of Γ_{jet} and θ_{LOS} . The hatched area indicates the limit set by 10 Schwarzschild radii for the $\sim 10^9 M_\odot$ central black hole. The dashed-dotted line shows the (arbitrary) injection radius $r_0 = 0''.5$ or 80 pc.















